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Hierarchical Neutrinos and Supersymmetric Inflation

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Abstract

A moderate extension of MSSM based on a left-right symmetric gauge group, within which hybrid inflation is 'naturally' realized, is discussed. The μ problem is solved via a Peccei-Quinn symmetry. Light neutrinos acquire hierarchical masses by the seesaw mechanism. They are taken from the small angle MSW resolution of the solar neutrino puzzle and the SuperKamiokande data. The range of parameters consistent with maximal $\nu_\mu - \nu_\tau$ mixing and the gravitino constraint is determined. The baryon asymmetry of the universe is generated through a primordial leptogenesis. The subrange of parameters, where the baryogenesis constraint is also met, is specified. The required values of parameters are more or less 'natural'.

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It is by now clear that the minimal supersymmetric standard model (MSSM), being unable to provide answers to a number of important physical issues, must be part of a more fundamental theory. Some of the shortcomings of MSSM are the following: i) Inflation cannot be implemented in its context. ii) There is no understanding of how the supersymmetric μ term, with $\mu \sim 10^2 - 10^3$ GeV, arises. iii) Neutrinos remain massless and, consequently, there are no neutrino oscillations in contrast to recent experimental evidence [1]. iv) Last, but not least, the observed baryon asymmetry of the universe cannot be easily generated (through the electroweak sphaleron processes).

Moderate extensions of MSSM provide [2–5] a suitable framework within which the hybrid inflationary scenario [6] is 'naturally' implemented. This means that a) there is no need for tiny coupling constants, b) the superpotential can have the most general form allowed by symmetry, c) supersymmetry guarantees that radiative corrections do not invalidate inflation, but rather provide the necessary slope along the inflationary trajectory, and d) supergravity corrections can be brought under control.

The μ problem, in these inflationary extensions of MSSM, could be solved [7] by coupling the inflaton system to the electroweak higgs superfields. In this case, however, the inflaton predominantly decays, after the end of inflation, into (s)higgses and the gravitino constraint [8] on the reheat temperature restricts [5,9] the relevant dimensionless coupling constants to 'unnaturally' small values ($\sim 10^{-5}$). Here, we will adopt an alternative solution [10,11] to the μ problem that relies on coupling the electroweak higgses to superfields causing the breaking of the Peccei-Quinn [12] symmetry.

Light neutrino masses can be generated either through the well-known seesaw mechanism after including right handed neutrino superfields or via coupling [5,11,13] the light neutrinos directly to $SU(2)_L$ triplet pairs of superfields of intermediate scale masses and 'tiny' vacuum expectation values (vevs). The latter possibility, studied in Ref. [5], is certainly more appropriate if light neutrinos are to provide the hot dark matter needed to explain [14] the large scale structure of the universe. Note that, for zero cosmological constant and almost flat spectrum of density perturbations, the hot dark matter density must be about 20% of the critical density keeping light neutrino masses in the eV range. Within a three neutrino scheme, compatibility with the atmospheric [1] and solar neutrino oscillations requires almost degenerate light neutrino masses which cannot be

'naturally' obtained from the seesaw mechanism.

Measurements [15] of the cluster baryon fraction combined with the low deuterium abundance constraint [16] on the baryon asymmetry suggest that the matter density is around 40% of the critical density of the universe ($\Omega_m \approx 0.4$). Also, recent observations [17] favor the existence of a cosmological constant whose contribution to the energy density can be as large as 60% of the critical density ($\Omega_\Lambda \approx 0.6$) driving the total energy density close to its critical value as required by inflation. The assumption that dark matter contains only a cold component leads then to a 'good' fit [18] of the cosmic background radiation and both the large scale structure and age of the universe data. Moreover, the possibility of improving this fit by adding light neutrinos as hot dark matter appears [19] to be rather limited and, thus, hierarchical neutrino masses are acceptable. Note that neutrino masses could be hierarchical even in the case of zero cosmological constant provided that hot dark matter consists of some other particles (say axinos).

Here, we will concentrate on the case of hierarchical neutrino masses generated by the seesaw mechanism. We consider a particular extension of MSSM based on a left-right symmetric gauge group within which hybrid inflation is 'naturally' implemented and the μ problem is solved via a Peccei-Quinn symmetry. Right handed neutrinos are present and form $SU(2)_R$ doublets with the right handed charged leptons. The inflaton decays into right handed neutrino superfields, whose subsequent decay produces [20] a primordial lepton asymmetry that is later converted into baryon asymmetry by electroweak sphaleron effects.

Our model is based on the left-right symmetric gauge group $G_{LR} = SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. The breaking of $SU(2)_R \times U(1)_{B-L}$ to $U(1)_Y$, at a superheavy scale $M \sim 10^{16}$ GeV, is achieved through the superpotential

$$W = \kappa S(l^c \bar{l}^c - M^2) \ , \tag{1}$$

where l^c, $\bar{l}^c$ is a conjugate pair of $SU(2)_R$ doublet left handed superfields with $B - L$ charges equal to 1, -1 respectively, and S is a gauge singlet left handed superfield. The coupling constant κ and the mass parameter M can be made real and positive by suitable phase redefinitions. The supersymmetric minima of the scalar potential lie on the D-flat direction $l^c = \bar{l}^{c*}$ at $\langle S \rangle = 0$, $|\langle l^c \rangle| = |\langle \bar{l}^c \rangle| = M$.

In this supersymmetric scheme, hybrid inflation [6] is automatically and 'naturally' realized [2–4,7,9]. The scalar potential possesses a built-in inflationary trajectory at $l^c = \bar{l}^c = 0$, $|S| > M$ with a constant tree-level potential energy density $\kappa^2 M^4$ which is responsible for the exponential expansion of the universe. Moreover, since this constant energy density breaks supersymmetry, there are important radiative corrections [3] which provide a slope along the inflationary trajectory necessary for driving the inflaton towards the vacua. At one-loop, the cosmic microwave quadrupole anisotropy is given [3,4] by

$$\left(\frac{\delta T}{T}\right)_Q \approx 4\sqrt{2}\pi \left(\frac{N_Q}{45}\right)^{1/2} \left(\frac{M}{M_P}\right)^2 x_Q^{-1} y_Q^{-1} \Lambda(x_Q^2)^{-1} \ , \tag{2}$$

where $N_Q \approx 50-60$ denotes the number of e-foldings experienced by our present horizon size during inflation, $M_P = 1.22 \times 10^{19}$ GeV is the Planck scale and

$$\Lambda(z) = (z-1)\ln(1-z^{-1}) + (z+1)\ln(1+z^{-1}) \ . \tag{3}$$

Also,

$$y_Q^2 = \int_1^{x_Q^2} \frac{dz}{z} \Lambda(z)^{-1} \ , \ y_Q \geq 0 \ , \tag{4}$$

with $x_Q = |S_Q|/M$ ($x_Q \geq 1$), S_Q being the value of the scalar field S when the scale which evolved to the present horizon size crossed outside the inflationary horizon. The superpotential parameter κ can be evaluated [3,4] from

$$\kappa \approx \frac{4\sqrt{2}\pi^{3/2}}{\sqrt{N_Q}} \ y_Q \ \frac{M}{M_P} \quad . \tag{5}$$

The μ term can be generated [7] by adding the superpotential coupling SH^2, where the electroweak higgs superfield $H = (H^{(1)}, H^{(2)})$ belongs to a bidoublet $(2,2)_0$ representation of $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. After gravity-mediated supersymmetry breaking, S develops [7] a vev $\langle S \rangle \approx -m_{3/2}/\kappa$, where $m_{3/2} \sim (0.1-1)$ TeV is the gravitino mass, and generates a μ term with $\mu = \lambda\langle S \rangle \approx -(\lambda/\kappa)m_{3/2}$.

This particular solution of the μ problem is [5,9], however, not totally satisfactory since it requires the presence of 'unnaturally' small coupling constants ($\kappa \lesssim \times 10^{-5}$). This is due to the fact that the inflaton system decays predominantly into electroweak higgs superfields via the renormalizable superpotential coupling SH^2. The gravitino constraint

[8] on the reheat temperature then severely restricts the corresponding dimensionless coupling constant and, consequently, the parameter κ. Moreover, for hierarchical neutrino masses from the seesaw mechanism, the requirement of maximal $\nu_\mu - \nu_\tau$ mixing from the SuperKamiokande experiment [1] further reduces [9] κ to become of order 10^{-6}.

An alternative solution of the μ problem will be adopted here. It is constructed [10] by coupling the electroweak higgses to superfields causing the breaking of the Peccei-Quinn [12] symmetry ($U(1)_{PQ}$) which solves the strong CP problem. For this purpose, we need two extra gauge singlet left handed superfields N and $\bar{N}$ with PQ charges -1 and 1. The relevant superpotential couplings are $N^2\bar{N}^2$ and N^2H^2. The scalar potential which is generated by the term $N^2\bar{N}^2$ after gravity-mediated supersymmetry breaking has been studied in Ref. [11]. For a suitable choice of parameters, the minimum lies at

$$|\langle N\rangle| \ = \ |\langle\bar{N}\rangle| \ \sim (m_{3/2}m_P)^{1/2} \ \sim 10^{11}\mathrm{GeV} \ , \tag{6}$$

where $m_P = M_P/\sqrt{8\pi}$ is the 'reduced' Planck mass. This scale is identified with the symmetry breaking scale f_a of $U(1)_{PQ}$. Substitution of $\langle N\rangle$ in the superpotential coupling N^2H^2 generates a μ parameter of order $m_{3/2}$ as desired.

This approach avoids the direct coupling of the inflaton to the electroweak higgses. Thus, the inflaton decays preferably to right handed neutrino superfields via non-renormalizable interactions, which are 'naturally' suppressed by m_P^{-1} (see below), rather than to higgses through renormalizable couplings. The gravitino constraint can then be satisfied for more 'natural' values of the dimensionless coupling constants.

The superpotential W of the full model contains, in addition to the terms in Eq.(1), the following couplings:

$$HQQ^c, \ HLL^c, \ N^2H^2, \ N^2\bar{N}^2, \ \bar{l}^c\bar{l}^cL^cL^c \ . \tag{7}$$

Here Q_i and L_i are the $SU(2)_L$ doublet left handed quark and lepton superfields, whereas Q_i^c and L_i^c are the $SU(2)_R$ doublet antiquarks and antileptons (i=1,2,3 is the family index). The quartic terms in Eq.(7) carry a factor m_P^{-1} which has been left out. Also, the dimensionless coupling constants as well as the family indices are suppressed. The last (non-renormalizable) term in Eq.(7) generates the intermediate scale Majorana masses of the right handed neutrinos after $SU(2)_R$ breaking.

The continuous global symmetries of this superpotential are $U(1)_B$ (and, consequently, $U(1)_L$) with the extra superfields S, l^c, $\bar{l}^c$, N, $\bar{N}$ carrying zero baryon number, an anomalous Peccei-Quinn symmetry $U(1)_{PQ}$, and a non-anomalous R-symmetry $U(1)_R$. The PQ and R charges of the superfields are as follows (W carries one unit of R charge):

$$PQ: \; H(1),\; Q(-1),\; Q^c(0),\; L(-1),\; L^c(0),\; S(0),\; l^c(0),\; \bar{l}^c(0),\; N(-1),\; \bar{N}(1) \;;$$

$$R: \; H(0),\; Q(1/2),\; Q^c(1/2),\; L(1/2),\; L^c(1/2),\; S(1),\; l^c(0),\; \bar{l}^c(0),\; N(1/2),\; \bar{N}(0) \;. \quad (8)$$

Note that $U(1)_B$ (and, thus, $U(1)_L$) invariance is automatically implied by $U(1)_R$ even if all possible non-renormalizable terms are included in the superpotential. This is due to the fact that the R charges of the products of any three color (anti)triplets exceed unity and cannot be compensated since there are no negative R charges available.

After $U(1)_L$ breaking by the vevs of l^c, $\bar{l}^c$, some lepton number violating effective operators will appear. In particular, the last term in Eq.(7) will generate the desirable intermediate scale masses for the right handed neutrinos. However, undesirable mixing of the higgs $H^{(2)}$ with L 's will also emerge from the allowed superpotential couplings $N\bar{N}LHl^c$ after the breaking of $U(1)_{PQ}$ by the vevs of N, $\bar{N}$. Also, the L 's will mix with $\bar{l}^c$ via the allowed couplings $N\bar{N}L^c\bar{l}^c$. To avoid such complications, we impose an extra discrete Z_2 symmetry ('lepton parity') under which L, L^c change sign.

The only superpotential terms which are permitted by the global symmetries $U(1)_R$, $U(1)_{PQ}$ and 'lepton parity' are the ones in Eqs.(1) and (7) as well as $LLl^cl^c\bar{N}^2l^c\bar{l}^c$ and LLl^cl^cHH modulo arbitrary multiplications by non-negative powers of the combination $l^c\bar{l}^c$. The vevs $\langle N\rangle$, $\langle\bar{N}\rangle$ together break $U(1)_{PQ}\times U(1)_R$ completely. Note that $U(1)_L$ is broken completely together with the gauge $U(1)_{B-L}$ by the superheavy vevs of l^c, $\bar{l}^c$. Thus, only $U(1)_B$ and 'lepton parity' remain exact.

As indicated above, after lepton number violation, the last term in Eq.(7) generates intermediate scale masses for the right handed neutrino superfields ν^c_i (i=1,2,3). The dimensionless coupling constant matrix of this term can be made diagonal with positive entries γ_i (i=1,2,3) by a rotation on ν^c_i 's. The right handed neutrino mass eigenvalues are then $M_i = 2\gamma_i M^2/m_P$ (with $\langle l^c\rangle$, $\langle\bar{l}^c\rangle$ taken positive by a $B-L$ transformation).

The light neutrino masses are generated via the seesaw mechanism in our scheme and, therefore, cannot be 'naturally' degenerate. We will, thus, assume hierarchical light

neutrino masses. Analysis [21] of the CHOOZ experiment [22] shows that the oscillations of solar and atmospheric neutrinos decouple. This fact allows us to concentrate on the two heaviest families ignoring the first one. We will denote the two positive eigenvalues of the light neutrino mass matrix by m_2 $(=m_{\nu_\mu})$, m_3 $(=m_{\nu_\tau})$. We take $m_{\nu_\mu} \approx 2.6 \times 10^{-3}$ eV which is the central value of the μ-neutrino mass coming from the small angle MSW resolution of the solar neutrino problem [23]. The τ-neutrino mass is taken to be $m_{\nu_\tau} \approx 7 \times 10^{-2}$ eV which is the central value implied by SuperKamiokande [1].

The determinant and the trace invariance of the light neutrino mass matrix imply [24] two constraints on the (asymptotic) parameters which take the form:

$$m_2 m_3 \ = \ \frac{\left(m_2^D m_3^D\right)^2}{M_2 \ M_3} \ , \tag{9}$$

$$m_2{}^2 + m_3{}^2 \ = \frac{\left(m_2^{D\ 2}\mathrm{c}^2 + m_3^{D\ 2}\mathrm{s}^2\right)^2}{M_2{}^2} +$$

$$\frac{\left(m_3^{D\ 2}\mathrm{c}^2 + m_2^{D\ 2}\mathrm{s}^2\right)^2}{M_3{}^2} + \frac{2(m_3^{D\ 2} - m_2^{D\ 2})^2\mathrm{c}^2\mathrm{s}^2\cos 2\delta}{M_2\, M_3} \ . \tag{10}$$

Here, $\mathrm{c} = \cos\theta, \ \mathrm{s} = \sin\theta$, and θ and δ are the rotation angle and phase which diagonalize the Majorana mass matrix of the right handed neutrinos in the basis where the 'Dirac' mass matrix of the neutrinos is diagonal with (positive) eigenvalues m_2^D, m_3^D $(m_2^D \leq m_3^D)$.

The $\mu - \tau$ mixing angle $\theta_{\mu\tau}$ lies [24] in the range

$$|\,\varphi - \theta^D| \leq \theta_{\mu\tau} \leq \varphi + \theta^D, \ \text{for} \ \varphi + \theta^D \leq \ \pi/2 \ , \tag{11}$$

where φ is the rotation angle which diagonalizes the light neutrino mass matrix in the basis where the 'Dirac' mass matrix is diagonal and θ^D is the 'Dirac' mixing angle (i.e., the 'unphysical' mixing angle with zero Majorana masses for the right handed neutrinos).

We now turn to the discussion of the inflaton decay. Here the inflaton consists [4] of the two complex scalar fields S and $\theta = (\delta\phi + \delta\bar{\phi})/\sqrt{2}$ $(\delta\phi = \phi - M, \ \delta\bar{\phi} = \bar{\phi} - M$ with ϕ, $\bar{\phi}$ being the neutral components of l^c, $\bar{l}^c)$ with a common mass $m_{infl} = \sqrt{2}\kappa M$. The scalar S (θ) can decay into a pair of bosonic (fermionic) ν_i^c 's via the last coupling in Eq.(7) and the coupling $\kappa S l^c \bar{l}^c$. The decay width is the same for both scalars and equals

$$\Gamma = \frac{1}{8\pi} \left(\frac{M_i}{M}\right)^2 m_{infl} \ . \tag{12}$$

Of course, decay of the inflaton into ν_i^c is only possible if $M_i < m_{infl}/2$. The reheat temperature is then given [4] by

$$T_r \ \approx \frac{1}{7} \left(\Gamma M_P\right)^{1/2} \ , \tag{13}$$

for MSSM spectrum, and the gravitino constraint [8] ($T_r \lesssim 10^9$ GeV) implies strong bounds on the M_i 's which satisfy the inequality $M_i < m_{infl}/2$. Consequently, the corresponding dimensionless coupling constants, γ_i , are restricted to be quite small.

To minimize the number of small couplings, we take $M_2 < m_{infl}/2 \leq M_3$ so that the inflaton decays into only one (the second heaviest) right handed neutrino with mass M_2. Moreover, we take $\gamma_3 = 1$, which gives $M_3 = 2M^2/m_P$. Using Eq.(5), the requirement $m_{infl}/2 \leq M_3$ becomes $y_Q \leq \sqrt{2N_Q}/\pi \approx 3.34$, for $N_Q = 55$, and Eq.(4) gives $x_Q \lesssim 3.5$. Eqs.(2) and (5) with $(\delta T/T)_Q \approx 6.6 \times 10^{-6}$ from COBE [25] are then used to calculate M and κ for each value of x_Q in this range. Eliminating x_Q, we obtain M as a function of κ depicted in Fig.1. The inflaton mass m_{infl} and the heaviest right handed neutrino mass M_3 are readily evaluated. The mass of the second heaviest right handed neutrino M_2 is restricted by the gravitino constraint [8] on T_r. We take it to be equal to its maximal allowed value in order to maximize γ_2. The value of M_2 is also depicted in Fig.1.

In our model, baryon number is conserved up to 'tiny' non-perturbative $SU(2)_L$ instanton effects. So the observed baryon asymmetry of the universe can only be produced by first generating a primordial lepton asymmetry [20] which is then partially converted into baryon asymmetry by non-perturbative electroweak sphaleron effects. The lepton asymmetry is produced through the decay of the superfield ν_2^c which emerges as decay product of the inflaton. This mechanism for leptogenesis has been discussed in Ref. [20]. The ν_2^c superfield decays into electroweak higgs and (anti) lepton superfields. The relevant one-loop diagrams are both of the vertex and self-energy type [26] with an exchange of ν_3^c. The resulting lepton asymmetry is [24]

$$\frac{n_L}{s} \approx 1.33 \ \frac{9T_r}{16\pi m_{infl}} \ \frac{M_2}{M_3} \ \frac{\mathrm{c}^2\mathrm{s}^2 \ \sin 2\delta \ (m_3^{D\,2} - m_2^{D\,2})^2}{|\langle H^{(1)}\rangle|^2 \ (m_3^{D\,2}\ \mathrm{s}^2 \ + \ m_2^{D\,2}\ \mathrm{c}^2)} \ , \tag{14}$$

where $|\langle H^{(1)}\rangle| \approx 174$ GeV. Note that this formula holds [27] provided that $M_2 \ll M_3$ and the decay width of ν_3^c is much smaller than $(M_3^2 - M_2^2)/M_2$, and both conditions are

well satisfied here. For MSSM spectrum, the observed baryon asymmetry is given [28] by $n_B/s = -(28/79)(n_L/s)$. It is important to ensure that the primordial lepton asymmetry is not erased by lepton number violating $2 \to 2$ scattering processes at all temperatures between T_r and 100 GeV. This gives [28] $m_{\nu_\tau} \lesssim 10$ eV which is readily satisfied.

For definiteness, we assume that the $\nu_\mu - \nu_\tau$ mixing is about maximal ($\theta_{\mu\tau} \approx \pi/4$) in accordance to the recent SuperKamiokande data [1]. We will also make the plausible assumption that the 'Dirac' mixing angle θ^D is negligible ($\theta^D \approx 0$). Under these circumstances, the rotation angle $\varphi \approx \pi/4$. Using the 'determinant' and 'trace' constraints in Eqs.(9) and (10) and diagonalizing the light neutrino mass matrix, we can determine the range of m_3^D which allows maximal $\nu_\mu - \nu_\tau$ mixing for each value of κ. These ranges are depicted in Fig.2 for all relevant values of κ and constitute the area in the $\kappa - m_3^D$ plane consistent with maximal mixing. For each allowed pair κ, m_3^D, the value of the phase δ leading to maximal mixing can be determined from the 'trace' condition. The corresponding lepton asymmetry is then found from Eq.(14). The line consistent with the low deuterium abundance constraint [16] on the baryon asymmetry of the universe ($\Omega_B h^2 \approx 0.019$) is also depicted in Fig.2. We see that the required values of κ ($\lesssim 4.2 \times 10^{-4}$), although somewhat small, are much more 'natural' than the ones encountered in previous models [9] that solved the μ problem and achieved maximal $\nu_\mu - \nu_\tau$ mixing. For these values of κ, the parameter $\gamma_2 \gtrsim 1.9 \times 10^{-3}$, which is quite satisfactory.

We have presented a moderate extension of MSSM based on a left-right symmetric gauge group. Hybrid inflation is 'naturally' realized and the μ problem is solved via a Peccei-Quinn symmetry. Light neutrinos acquire hierarchical masses by the seesaw mechanism, which are taken to be consistent with the small angle MSW resolution of the solar neutrino problem and the SuperKamiokande data. Under these conditions, we determine the range of κ and the (asymptotic) 'Dirac' neutrino masses consistent with maximal $\nu_\mu - \nu_\tau$ mixing and the gravitino constraint. The baryon asymmetry of the universe is generated through a primordial leptogenesis. We specify the subrange of parameters where the baryogenesis constraint is also met. The required values of the relevant coupling constants turn out to be more or less 'natural'.

REFERENCES

[1] T. Kajiata, talk given at the XVIIIth International Conference on Neutrino Physics and Astrophysics (Neutrino'98), Takayama, Japan, 4-9 June, 1998.

[2] E. J. Copeland, A. R. Liddle, D. H. Lyth, E. D. Stewart and D. Wands, Phys. Rev. **D49** (1994) 6410.

[3] G. Dvali, Q. Shafi and R. Schaefer, Phys. Rev. Lett. **73** (1994) 1886.

[4] G. Lazarides, R. Schaefer and Q. Shafi, Phys. Rev. **D56** (1997) 1324.

[5] G. Lazarides, hep-ph/9812454 (to appear in Phys. Lett. B).

[6] A. D. Linde, Phys. Lett. **B259** (1991) 38; Phys. Rev. **D49** (1994) 748.

[7] G. Dvali, G. Lazarides and Q. Shafi, Phys. Lett. **B424** (1998) 259.

[8] M. Yu. Khlopov and A. D. Linde, Phys. Lett. **138B** (1984) 265; J. Ellis, J. E. Kim and D. Nanopoulos, Phys. Lett. **145B** (1984) 181.

[9] G. Lazarides and N. D. Vlachos, Phys. Lett. **B441** (1998) 46.

[10] J. E. Kim and H. P. Nilles, Phys. Lett. **138B** (1984) 150.

[11] G. Lazarides and Q. Shafi, Phys. Rev. **D58** (1998) 071702.

[12] R. Peccei and H. Quinn, Phys. Rev. Lett. **38** (1977) 1440; S. Weinberg, Phys. Rev. Lett. **40** (1978) 223; F. Wilczek, Phys. Rev. Lett. **40** (1978) 279.

[13] G. Lazarides, Q. Shafi and C. Wetterich, Nucl. Phys. **B181** (1981) 287; C. Wetterich, Nucl. Phys. **B187** (1981) 343; R. N. Mohapatra and G. Senjanovic, Phys. Rev. **D23** (1981) 165; R. Holman, G. Lazarides and Q. Shafi, Phys. Rev. **D27** (1983) 995.

[14] Q. Shafi and F. W. Stecker, Phys. Rev. Lett. **53** (1984) 1292. For a recent review and other references see Q. Shafi and R. K. Schaefer, hep-ph/9612478.

[15] A. E. Evrard, MNRAS **292** (1997) 289.

[16] S. Burles and D. Tytler, Astrophys. Journal **499** (1998) 699; ibid. **507** (1998) 732.

[17] B. Schmidt *et al.*, Astrophys. Journal **507** (1998) 46; A. Riess *et al.*, Astron. J. **116** (1998) 1009; S. Perlmutter *et al.*, astro-ph/9812473; S. Perlmutter *et al.*, astro-ph/9812133 (to appear in Astrophys. J.).

[18] L. Krauss and M. S. Turner, Gen. Rel. Grav. **27** (1995) 1137; J. P. Ostriker and P. J. Steinhardt, Nature **377** (1995) 600; A. R. Liddle *et al.*, MNRAS **282** (1996) 281.

[19] J. R. Primack and M. A. K. Gross, astro-ph/9810204.

[20] M. Fukugita and T. Yanagida, Phys. Lett. **B174** (1986) 45; W. Buchmüller and M. Plümacher, Phys. Lett. **B389** (1996) 73. In the context of inflation see G. Lazarides and Q. Shafi, Phys. Lett. **B258** (1991) 305; G. Lazarides, C. Panagiotakopoulos and Q. Shafi, Phys. Lett. **B315** (1993) 325, (E) ibid. **317** (1993) 661.

[21] C. Giunti, hep-ph/9802201.

[22] M. Apollonio *et al.*, Phys. Lett. **B420** (1998) 397.

[23] A. Smirnov, hep-ph/9611465 and references therein.

[24] G. Lazarides, Q. Shafi and N. D. Vlachos, Phys. Lett. **B427** (1998) 53.

[25] G. F. Smoot *et al.*, Astrophys. J. Lett. **396** (1992) L1; C. L. Bennett *et al.*, Astrophys. J. Lett. **464** (1996) 1.

[26] L. Covi, E. Roulet and F. Vissani, Phys. Lett. **B384** (1996) 169.

[27] A. Pilaftsis, Phys. Rev. **D56** (1997) 5431.

[28] L. E. Ibáñez and F. Quevedo, Phys. Lett. **B283** (1992) 261.

FIGURES

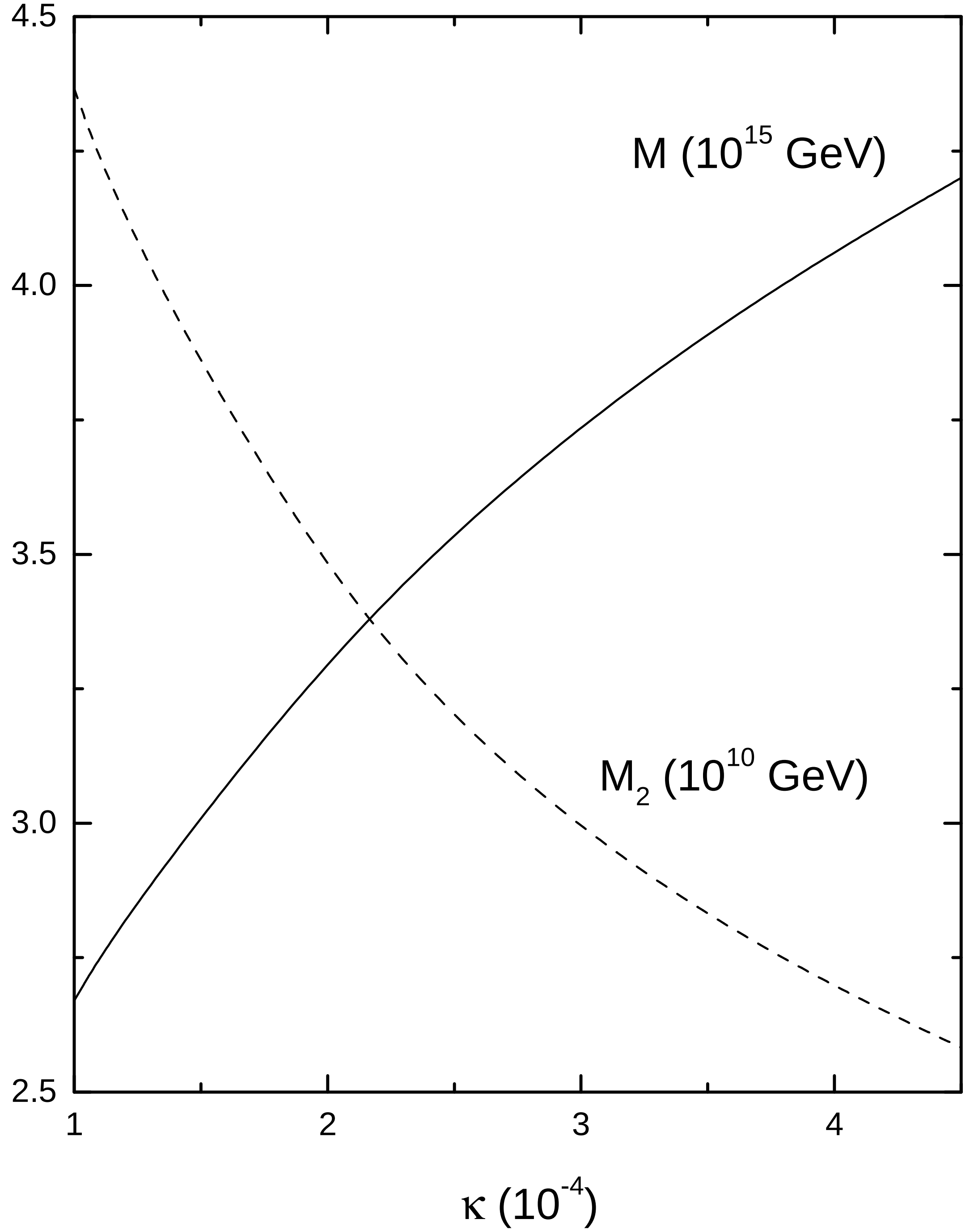

FIG. 1. The mass scale M (solid line) and the Majorana mass of the second heaviest right handed neutrino M_2 (dashed line) as functions of κ.

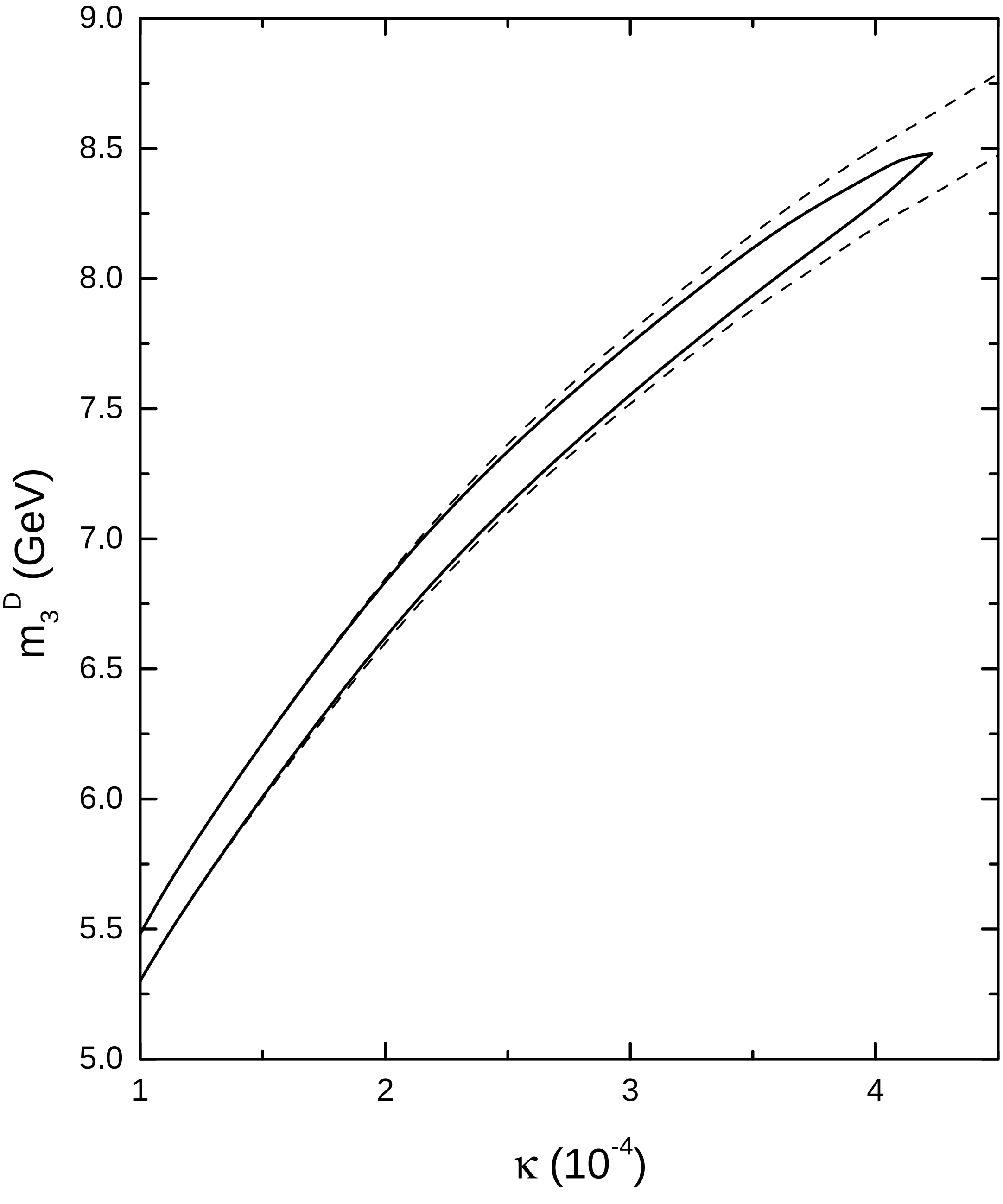

FIG. 2. The area (bounded by the dashed lines) on the $\kappa - m_3^D$ plane consistent with maximal $\nu_\mu - \nu_\tau$ mixing and the gravitino constraint. Along the thick solid line the low deuterium abundance constraint on the baryon asymmetry of the universe is also satisfied.